The Thomas Jefferson National Accelerator Facility
Theory Group Preprint Series

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The Southeastern Universities Research Association (SURA) operates the Thomas Jefferson National Accelerator Facility for the United States Department of Energy under contract DE-AC05-84ER40150.

JLAB-THY-99-03

SPECTROSCOPY – AN INTRODUCTION AND OVERVIEW

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The study of baryons can provide us with critical insights into the nature of QCD in the confinement domain. The key to progress in this field is the identification of its important degrees of freedom. I explain why I believe that the adiabatic approximation is central to understanding the absence of gluonic excitations at low energies, and describe an extension of this approximation which can help us to understand the resiliency of the valence quark model to meson loop corrections. I close with a survey of issues that I hope to see resolved before Baryon 2001.

1 Why Baryons?

There are, appropriately enough, three main reasons why I believe that baryons deserve the special attention they get from the series of conferences to which Baryons '98 belongs.

The first is that baryons *are* the stuff of which our world is made. As such they must be at center stage in any discussion of the nitty-gritty of why the world we actually experience has the character it does. Thus an understanding of how Quantum Chromodynamics (QCD) makes baryons must form the basis for an eventual understanding of the origin of the forces between nucleons and thence of the origin of the periodic table. I am convinced that completing this chapter in the history of science will be one of the most interesting and fruitful areas of physics for the next twenty years.

My second reason is that they are the simplest system in which the essential nonabelian character of QCD is manifest. There are, after all, N_c quarks in a proton because there are N_c colors, and this fact is in turn a consequence of the remarkable and quintessentially nonabelian property of QCD that three particles can attract each other (in contrast to Quantum Electrodynamics (QED) where the e^-e^- force is repulsive).

The third reason is historical. It is no accident that baryons have played a much more prominent role in the discovery of QCD than mesons. Gell-Mann and Zweig were forced to the quarks by $3 \times 3 \times 3$ giving the octet and decuplet; Greenberg was led to color by the spin-statistics paradox in the Δ^{++}, and Dalitz's quark model for baryons was one of the earliest indications of the power of the valence quark model.

1.1 *What is the Goal?*

What is the goal of this research, and indeed of all modern work on QCD? It should be to understand "strong QCD" [1] and nuclei. This includes being able to compute some quantities with precision, but most importantly acheiving a qualitative explanation of the main features of strong QCD, including the answers to such questions as:

- What is the physical origin of color confinement?
- Why is the low energy spectrum dominated by what appear to be $q\bar{q}$ and qqq systems?
- Where, as a corollary, are the excitations of the gluonic degrees of freedom?
- Where are the excitations of the sea quark degrees of freedom?
- Why is a nucleus made of nucleons, instead of a "quark soup"?
- What is the origin of the well-established empirical nucleon-nucleon force?

1.2 *What is the Key?*

I believe that the key to a qualitative understanding of strong QCD is the same as in most other areas of physics: identifying the appropriate degrees of freedom. For example, atomic physics is based on taking the nuclei and electrons as the low energy effective degrees of freedom, with the underlying effects of nucleons subsumed into static nuclear properties and those of photons into low energy effective potentials; nuclear physics is in turn very well-described by nucleons moving in an empirical nucleon-nucleon potential.

2 The Failure of Quarks and Gluons

We all know that asymptotic freedom guarantees that at sufficiently small distances QCD becomes a weakly coupled quark-gluon theory which is amenable to a perturbative expansion in the running coupling constant α_s. However, the other side of this coin is that at large distances α_s becomes large so that quark-gluon perturbation theory breaks down.

In fact, we now know from numerical studies that QCD predicts confinement: the potential energy between two static quarks grows linearly with their separation r with a constant of proportionality b, called the string tension, that

is about 1 GeV/fm. Let me show you that such a result rigorously implies the breakdown of perturbative QCD. Given that confinement is the central feature of strong QCD, we are therefore forced to seek new methods for the study of most strong interaction phenomena.

In the pure gluon sector of QCD in which the static potential problem is posed (*i.e.*, QCD with static sources and *no* dynamical quarks), the equation for the string tension must take the form

$$b = f_b(g^2) \tag{1}$$

where f_b is some function of the dimensionless coupling constant g^2 since this is the only parameter of pure QCD. This equation is impossible, however, since b has dimensions of $[mass]^2$! The resolution of this paradox lies in the fact that g^2 is not a coupling "constant": according to asymptotic freedom

$$\frac{1}{g^2(Q^2)} = \frac{1}{g_0^2(Q_0^2)} + \frac{11}{16\pi^2}\ln\frac{Q^2}{Q_0^2} \tag{2}$$

where $g(g_0)$ is the effective coupling at momentum transfer $Q^2(Q_0^2)$. Thus QCD is defined by a universal "coupling constant curve" $g^2(Q^2)$ on which g^2 takes all values from zero to infinity, and not a single number. In a given universe with scales external to QCD (like the electroweak electron mass or the masses of the current quarks) this universal curve can be "pegged" to a given normalization at some external scale μ^2, but in pure QCD this is irrelevant: for us the key point is that a particular curve can be defined by choosing a value for $g^2(\mu^2)$ at any normalization point μ^2. This choice then simultaneously gives us a coupling constant $g^2(\mu^2)$ and a scale to give dimension to equation (1):

$$b = \mu^2 f_b(g^2(\mu^2)) \ . \tag{3}$$

Thus in a pure QCD world, the string tension b and all other dimensionful quantities would have a scale set by the dummy variable μ^2, and all observables would be dimensionless ratios in which this variable cancels out.

Now note that any point μ^2 could have been chosen to define the curve $g^2(Q^2)$ and so

$$\frac{db}{d\mu^2} = 0 \tag{4}$$

or, *i.e.*,

$$0 = f_b - \frac{11}{16\pi^2} g^4 \frac{df_b}{dg^2} \tag{5}$$

implying that

$$f_b \propto exp\left[-\frac{16\pi^2}{11g^2}\right] \ . \tag{6}$$

This essential singularity in g^2 means that the "Feynman diagrammar" is useless for this problem, and that *plane wave quarks and gluons are not a useful starting point for low-energy, confinement-dominated physics.* To make progress in understanding the main phenomena of strong interaction physics, we must therefore either resort to purely numerical methods (*e.g.*, lattice QCD), or we must replace the Feynman diagrammar by new conceptual elements.

3 All Roads Lead to Valence Quarks and Flux Tubes

From the preceeding discussion on the inappropriateness of the perturbative quark and gluon degrees of freedom for describing the phenomena of strong QCD, it will come as no surprise that foremost among the puzzles we face is in fact a glaring "degree of freedom" problem: the established low energy spectrum of QCD behaves as though it is built from the degrees of freedom of spin-$\frac{1}{2}$ fermions confined to a $q\bar{q}$ or qqq system. Thus, for mesons we seem to observe a "quarkonium" spectrum, while for the baryons we seem to observe the spectrum of the two relative coordinates of three spin-$\frac{1}{2}$ degrees of freedom.

These apparent degrees of freedom are to be contrasted with the most naïve interpretation of QCD in terms of the perturbative degrees of freedom which would lead us to expect a low energy spectrum exhibiting 36 quark and antiquark degrees of freedom (3 flavors $\times$ 2 spins $\times$ 3 colors for particle and antiparticle), *and* 16 gluon degrees of freedom (2 spins $\times$ 8 colors).

The second major "degree of freedom problem" has to do with $q\bar{q}$ pair creation. At least naïvely, one would expect pair creation to be so strong that a valence quark model would fail dramatically. That pair creation should be expected to lead to dramatic failures of valence quark model spectroscopy and dynamics is true even though we know empirically (and understand theoretically) that pair creation is suppressed, *i.e.*, that the observed hadronic spectrum is dominated by relatively narrow resonances. I will discuss this problem and how it might be resolved in the next Section.

First let me describe the compelling insights coming from three different directions which converge on a simple valence quark plus glue picture of the structure of strong QCD.

3.1 The Large N_c Limit of QCD

It is now widely appreciated that many of the observed features of the strong interactions can be rationalized in QCD within the $1/N_c$ expansion [2]. Moreover, there is growing evidence from lattice QCD that while $N_c = 3$ might not be sufficiently large for the $1/N_c$ expansion to be used quantitatively, the

main qualitative features of QCD (including confinement and the spontaneous breakdown of chiral symmetry) are independent of N_c.

We should therefore take seriously the fact that it can be shown in the large-N_c limit that hadron two-point functions are dominated by graphs in which the valence quark lines propagate from their point of creation to their point of annihilation without additional quark loops. A form of the OZI rule [3] also emerges naturally. Large-N_c QCD thus presents a picture of narrow resonances interacting weakly with hadronic continua. In this picture the resonances themselves are made of valence quarks and glue.

3.2 Quenched QCD

Quenched lattice QCD provides other new insights into QCD. In quenched QCD the lattice sums amplitudes over all time histories in which no $q\bar{q}$ loops are present. It thus gives quantitative results from an approximation with many elements in common with the large N_c limit. One of the most remarkable features of these calculations is that despite what would seem to be a drastic approximation, they provide a reasonably good description of low energy phenomenology. Indeed, for various intermediate quantities like the QCD string tension they provide very good approximations to full QCD results with the true lattice coupling constant replaced by an effective one. In quenched QCD, as in the large N_c limit, two point functions thus seem to be well-approximated by their valence content (namely pure glue for glueballs, $q\bar{q}$ plus glue for mesons, and qqq plus glue for baryons).

In comparing the large N_c limit and quenched lattice QCD we note that:

- In both pictures all resonances have only valencc quarks, but they have an unlimited number of gluons. Thus they support valence models for mesons and baryons, but not for glueballs or for the gluonic content of mesons and baryons.

- In both pictures a propagating valence quark has contributions from not only a positive energy quark propagator, but also from "Z-graphs". (A "Z-graph" is a time-ordered graph in which the interactions first produce a pair and then annihilate the antiparticle of the produced pair against the original propagating particle). Cutting through a two-point function at a fixed time therefore would in general reveal not only the valence quarks but also a large $q\bar{q}$ sea. This does not seem to correspond to the usual valence approximation. Consider, however, the Dirac equation for a single light quark interacting with a static color source (or a single light quark confined in a bag). This equation represents the sum of a set of Feynman graphs which also include Z-graphs, but the effects of those graphs is captured in the lower components of the single particle Dirac spinor. *I.e.*, such Z-graphs correspond to relativistic corrections to the quark model. That such corrections are important in the quark model has been known for a long time. For us the important point is that while they have quantitative effects on quark model predictions (*e.g.*, they are commonly held to be responsible for much of the required reduction of the nonrelativistic quark model prediction that $g_A = 5/3$ in neutron beta decay), they do not qualitatively change the single-particle nature of the spectrum of the quark of our example, nor would they qualitatively change the spectrum of $q\bar{q}$ or qqq systems.

- Finally, we note that the large N_c and quenched approximations are *not* identical. For example, the NN interaction is a $1/N_c$ effect, but it is not apparently suppressed in the quenched approximation.

3.3 The Heavy Quark Limit

The third perspective from which there is support for the same picture is the heavy quark limit[4]. While this limit has the weakest theoretical connections to the light quark world, it has powerful phenomenological connections: see Fig. 1. We see from this picture that in mesons containing a single heavy quark, $\Delta E_{orbital}$ (the gap between, for example, the $J^{PC} = 1^{--}$ and 2^{++} states), is approximately independent of m_Q, as predicted in the heavy quark limit, while $\Delta E_{hyperfine}$ varies like m_Q^{-1} as expected.

Recall that in the heavy quark limit a hadronic two-point function is dominated by a single valence Q plus its associated "brown muck", with neither $Q\bar{Q}$ loops nor Q Z-graphs. The fact that heavy-quark-like behaviour persists all the way down to light quark masses suggests that light quarks, like heavy quarks, behave like single valence quarks and thus by extension that the "brown muck" behaves like a single valence antiquark.

Fig. 2 shows that heavy quark behaviour also apparently persists in a stronger form: the light meson spectrum appears to mimic the $Q\bar{Q}$ quarkonium spectrum. This is surprising since this latter spectrum depends on the decoupling of gluonic excitations (as opposed to glue) from the spectrum *via* an adiabatic approximation.

In view of this surprise, let us examine the nature of the adiabatic approximation in QCD. Consider first QCD without dynamical quarks in the presence of fixed $Q_1\bar{Q}_2$ or $Q_1Q_2Q_3$ sources[5,6]. The ground state of QCD with these sources in place will be modified, as will be its excitation spectrum. For excitation energies below those required to produce a glueball, this spectrum

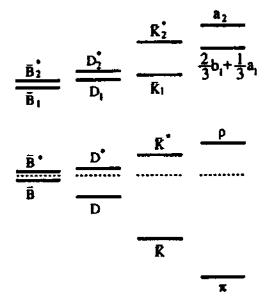

Figure 1: The $Q\bar{q}$ meson spectra as a function of the "heavy" quark mass.

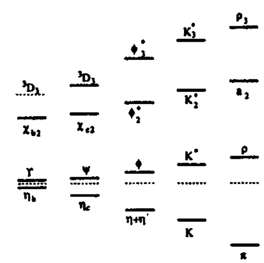

Figure 2: The $Q\bar{Q}$ meson spectra as a function of the "heavy" quark mass.

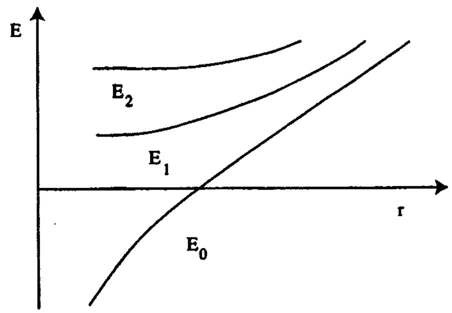

Figure 3: Schematic of the low-lying adiabatic surfaces of $Q_1\bar{Q}_2$ at separation r; $E_0(r)$ is the gluonic ground state, $E_1(r)$ the first excited state, *etc.*.

will presumably be discrete for each value of the $Q_1\bar{Q}_2$ relative spatial separation $\vec{r}$, with each eigenvalue being a continuous function of $|\vec{r}|$, as shown schematically in Fig. 3. There will be analogous spectra for $Q_1Q_2Q_3$ which are functions of its two relative coordinates. We call the energy surface traced out by a given level of excitation as the positions of the sources are varied an adiabatic surface, and *define* the "quark model limit" to be applicable when the quark sources move along the lowest adiabatic surface.

We all know a simple molecular physics analogy to this approximation. Diatomic molecular spectra can be described in an adiabatic approximation by holding the two relevant atomic nuclei at fixed separation r and then solving the Schrödinger problem for the (mutually interacting) electrons moving in the static electric field of the nuclei. The electrons will, for fixed r, have a ground state and excited states which will eventually become a continuum above energies required to ionize the molecule. The resulting adiabatic energy functions (when added to the internuclear Coulomb energy) then serve as effective internuclear potentials on which vibration-rotation spectra can be built. Molecular transitions can then take place within states built on a given surface or between surfaces.

In the "quark model limit" the quark sources play the rôle of the nuclei, and the glue plays the rôle of the electrons. From this point of view we can see

clearly that conventional meson and baryon spectroscopy has only scratched the surface of even $q_1\bar{q}_2$ and $q_1q_2q_3$ spectroscopy: so far we have only studied the vibration-rotation bands built on the lowest adiabatic surface corresponding to the gluonic ground state. We should expect to be able to build other "hadronic worlds" on the surfaces associated with excited gluonic states[6].

While the adiabatic approximation is more general, it is becoming increasingly firmly established that this approximation is realized in QCD in terms of the development of a confining chromoelectric flux tube. These flux tubes are the analog of the Abrikosov vortex lines that can develop in a superconductor subjected to a magnetic field, with the vacuum acting as a dual (*i.e.*, electric) superconductor creating a chromoelectric Meissner effect. A $Q\bar{Q}$ system held at fixed separation $r >> \Lambda_{QCD}$ is known to have as its ground state a flux tube which leads to an effective low energy (adiabatic) potential corresponding to the standard "quarkonium" potential. However, this system also has excited states, corresponding to gluonic adiabatic surfaces in which a phonon has been excited in the flux tube, and on which spectra of "hybrid states" are built.

Lattice results allow us to check many aspects of the flux tube picture. For example, the lattice confirms the flux tube model prediction that sources with triality are confined with a string tension proportional to the square of their color Casimir. The predicted strongly collimated chromoelectric flux lines have also been seen on the lattice. I have found it particularly encouraging that the first excited adiabatic surfaces have been seen[7] with an energy gap $\delta V(r) = \pi/r$ above the quarkonium potential as predicted[6], and with the expected doubly-degenerate phonon quantum numbers (see Fig. 4). This strongly suggests that the J^{PC} exotic hybrid mesons predicted ten years ago[6] exist.

The flux tube model thus offers a possible explanation for one of the most puzzling apparent inconsistencies between the naïve quark model and QCD. Moreover, as discussed above, in the large N_c limit of QCD hadrons do indeed consist of just their valence quarks and the glue between them. Thus the flux tube model may legitimately be viewed as a candidate realization of QCD in the large N_c limit which is in addition consistent with insights into strong QCD which have emerged from quenched lattice QCD and from heavy quark theory.

4 Unquenching the Quark Model

The valence quark plus glue dominance embodied in the flux tube model can at best be a starting point for a systematic description of strong QCD, since we know that $q\bar{q}$ pair creation plays an important rôle in many phenomena. Nevertheless, näively attempting to add $q\bar{q}$ pair creation to the valence quark

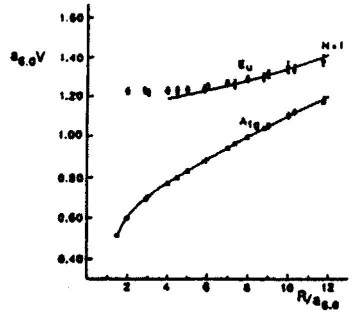

Figure 4: The ground state and first excited adiabatic potentials from lattice QCD [7].

model leads to a number of very serious problems. These problems and potential solutions to them have been extensively discussed in a series of papers on "unquenching" the quark model [8,9,10]. In the following I briefly describe these problems and solutions.

4.1 The Origin and Resiliency of Potential Models

In the preceeding Section, we saw that one of the "degrees of freedom" problems of the valence quark model could plausibly be solved by the flux tube model: the apparent absence of low energy degrees of freedom associated with the glue. The second "degree of freedom" problem is the apparent absence of excitations associated with the strong $q\bar{q}$ sea. Very closely related to this puzzle is the apparent unimportance of strong meson loop corrections.

A simple resolution of this puzzle has been proposed which is an extenion of the adiabatic approximation to the flux tube: the *full* quark potential model arises from an adiabatic approximation *including* the extra $q\bar{q}$ degrees of freedom embodied in the flux tube. At short distances where perturbation theory applies, the effect of N_f types of light $q\bar{q}$ pairs is (in lowest order) to shift the coefficient of the Coulombic potential from $\alpha_s^{(0)}(Q^2) = \frac{12\pi}{33 ln(Q^2/\Lambda_0^2)}$ to

$\alpha_s^{(N_f)}(Q^2) = \frac{12\pi}{(33-2N_f) ln(Q^2/\Lambda_{N_f}^2)}$. The net effect of such pairs is thus to produce a *new* effective short distance $Q\bar{Q}$ potential. Similarly, when pairs bubble up in the flux tube (*i.e.*, when the flux tube breaks to create a $Q\bar{q}$ plus $q\bar{Q}$ system and then "heals" back to $Q\bar{Q}$), their net effect is to cause a shift $\Delta E_{N_f}(r)$ in the ground state gluonic energy which in turn produces a new long-range effective $Q\bar{Q}$ potential.

It has indeed been shown [8] that the net long-distance effect of the bubbles is to create a new string tension b_{N_f} (*i.e.*, that the potential remains linear). Since this string tension is to be associated with the observed string tension, after renormalization *pair creation has no effect on the long-distance structure of the quark model in the adiabatic approximation.* Thus the net effect of mass shifts from pair creation is much smaller than one would naïvely expect from the typical width Γ: shifts that are not absorbed into the physical string tension can only arise from nonadiabatic effects. For heavy quarkonium, these shifts can in turn be associated with states which are strongly coupled to nearby thresholds [11].

It should be emphasized that no simple truncation of the set of all meson loop graphs can reproduce such results: to recover the adiabatic approximation requires summing over large towers of $Q\bar{q}$ plus $q\bar{Q}$ intermediate states to allow a duality with the $q\bar{q}$ loop diagrams which have strength at high energy.

4.2 The Survival of the OZI Rule

There is another puzzle of hadronic dynamics which is reminiscent of this one: the success of the OZI rule [3]. A generic OZI-violating amplitude A_{OZI} can also be shown to vanish like $1/N_c$. However, there are several unsatisfactory features of this "solution" to the OZI mixing problem [12]. Consider ω-ϕ mixing as an example. This mixing receives a contribution from the virtual hadronic loop process $\omega \to K\bar{K} \to \phi$, both steps of which are OZI-allowed, and each of which scales with N_c like $\Gamma^{1/2} \sim N_c^{-1/2}$. The large N_c result that this OZI-violating amplitude behaves like N_c^{-1} is thus not peculiar to large N_c: it just arises from "unitarity" in the sense that the real and imaginary parts of a generic hadronic loop diagram will have the same dependence on N_c. The usual interpretation of the OZI rule in this case --- that "double hairpin graphs" are dramatically suppressed --- is untenable in the light of these OZI-allowed loop diagrams, which expose the deficiency of the large N_c argument: $A_{OZI} \sim \Gamma$ is *not* a good representation of the OZI rule. (Continuing to use ω-ϕ mixing as an example, we note that $m_\omega - m_\phi$ is numerically comparable to a typical hadronic width, so the large N_c result would predict an ω-ϕ mixing angle of order unity in contrast to the observed pattern of very weak mixing

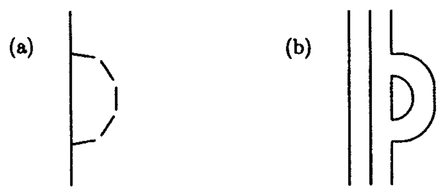

Figure 5: A meson loop correction to a baryon propagator, drawn at (a) the hadronic level, and (b) the quark level.

which implies that $A_{OZI} << \Gamma << m$.)

Unquenching the quark model thus endangers the naïve quark model's agreement with the OZI rule. It has been shown[9] how this disaster is naturally averted in the flux tube model through a "miraculous" set of cancellations between mesonic loop diagrams consisting of apparently unrelated sets of mesons (*e.g.*, the $K\bar{K}$, $K\bar{K}^* + K^*\bar{K}$, and $K^*\bar{K}^*$ loops tend to strongly cancel against loops containing a K or K^* plus one of the four strange mesons of the $L = 1$ meson nonets). Of course the "miracle" occurs for a good reason: the sum of *all* hadronic loops is dual to two $q\bar{q}$ hairpins of different flavors created and destroyed by a $^3P_0$ operator[13,14,15], but in the closure approximation such an operator cannot create mixing in other than a scalar channel.

While slightly more complex, it should be noted that OZI-violating baryon couplings like $p \to p\phi$ can be induced by loop diagrams which are essentially identical to those responsible for $\omega - \phi$ mixing. The same mechansism which prevents $\omega - \phi$ mixing also prevents such OZI-violating couplings.

4.3 Current Matrix Elements of the $q\bar{q}$ Sea

The preceeding perspective on the origin and dynamics of OZI-violation has an immediate impact on the issue of the $s\bar{s}$ content of the proton, and indeed more generally on the current matrix elements of the $q\bar{q}$ sea. Consider the $s\bar{s}$ content of the proton from the quark-level process shown in Fig. 5. The main new feature which this discussion forces upon us is a sum over a *complete set* of strange intermediate states, rather than just a few low-lying states[10].

The lower vertex in Fig. 5 arises when $s\bar{s}$ pair creation perturbs the initial nucleon state vector so that, to leading order in pair creation,

$$|p\rangle \to |p\rangle + \sum_{Y^*K^*\ell S} \int q^2 dq \; |Y^*K^*q\ell S\rangle \frac{\langle Y^*K^*q\ell S|h_{s\bar{s}}|p\rangle}{M_p - E_{Y^*} - E_{K^*}} , \qquad (7)$$

where $h_{s\bar{s}}$ is the $s\bar{s}$ quark pair creation operator, Y^* (K^*) is the intermediate baryon (meson), q and ℓ are the relative radial momentum and orbital angular momentum of Y^* and K^*, and S is the sum of their spins. This process will generate non-zero expectation values for strangeness observables:

$$\begin{aligned} \langle O_s \rangle \; = \; & \sum_{\substack{Y^*K^*\ell S \\ Y^{*\prime}K^{*\prime}\ell' S'}} \int q^2 dq \, q'^2 dq' \frac{\left\langle p|h_{s\bar{s}}|Y^{*\prime}K^{*\prime}q'\ell' S'\right\rangle}{M_p - E_{Y^{*\prime}} - E_{K^{*\prime}}} \\ & \times \left\langle Y^{*\prime}K^{*\prime}q'\ell' S'|O_s|Y^*K^*q\ell S\right\rangle \frac{\langle Y^*K^*q\ell S|h_{s\bar{s}}|p\rangle}{M_p - E_{Y^*} - E_{K^*}} . \qquad (8) \end{aligned}$$

The derivation of this simple equation, including the demonstration that it is gauge invariant, is straightforward[10]. Some cases of interest include $O_s = \Delta s$, R_s^2, and μ_s, where Δs is the quantity relevant to the proton spin crisis, and where R_s^2 and and μ_s are the strangeness radius and magnetic moment, respectively.

It is useful to refer to the closure-spectator limit of Eq. (8). This is the limit in which the energy denominators do not depend strongly on the quantum numbers of Y^* and K^*, so that the sums over intermediate states collapse to 1, giving

$$\langle O_s \rangle \propto \langle p|h_{s\bar{s}} O_s h_{s\bar{s}}|p\rangle \propto \langle 0|h_{s\bar{s}} O_s h_{s\bar{s}}|0\rangle \; , \qquad (9)$$

where the second step follows since $h_{s\bar{s}}$ does not couple to the motion of the valence spectator quarks. We see that the expectation value of O_s is taken between the $^3P_0$ states created by $h_{s\bar{s}}$. From the J^{PC} of the $^3P_0$ pair it then follows that $\Delta s = R_s^2 = \mu_s = 0$ in the closure-spectator limit (a result which would not be seen if only the lowest term, or lowest few terms, were included in the closure sum).

Strange Quark Contributions to the Proton Spin, Charge Radius, and Magnetic Moment

Δs, the fraction of the proton's spin carried by strange quarks, is given by twice the expectation value of the s and $\bar{s}$ spins :

$$\Delta s = 2\left\langle S_z^{(s)} + S_z^{(\bar{s})} \right\rangle . \qquad (10)$$

Let us first examine the contribution to Δs from just the lowest-lying intermediate state, ΛK. The P-wave ΛK state with $J = J_z = \frac{1}{2}$ is

$$\left|(\Lambda K)_{P\frac{1}{2}}\right\rangle = \sqrt{\frac{2}{3}}\left|(\Lambda_{\downarrow} K)_{m=1}\right\rangle - \sqrt{\frac{1}{3}}\left|(\Lambda_{\uparrow} K)_{m=0}\right\rangle. \tag{11}$$

The $\bar{s}$ quark in the kaon is unpolarized, while the s quark in the Λ carries all of the Λ's spin; because of the larger coefficient multiplying the first term in (11), the ΛK intermediate state alone gives a negative contribution to Δs.

When we add in the $(\Lambda K^*)_{P\frac{1}{2}}$ and $(\Lambda K^*)_{P\frac{3}{2}}$ states (note that the subscripts denote the quantities ℓS defined previously), we have

$$\Delta s \propto \begin{pmatrix} 1 & -\sqrt{\frac{1}{3}} & \sqrt{\frac{8}{3}} \end{pmatrix} \frac{1}{18} \begin{bmatrix} -3 & \sqrt{3} & -\sqrt{24} \\ & -1 & \sqrt{8} \\ & & 10 \end{bmatrix} \begin{pmatrix} 1 \\ -\sqrt{\frac{1}{3}} \\ \sqrt{\frac{8}{3}} \end{pmatrix} \tag{12}$$

in the closure limit. Here the matrix is just $2(S_z^{(s)} + S_z^{(\bar{s})})$ (which is of course symmetric), and the vectors give the relative coupling strengths of the proton to $[\ (\Lambda K)_{P\frac{1}{2}}, (\Lambda K^*)_{P\frac{1}{2}}, (\Lambda K^*)_{P\frac{3}{2}}\ ]$. Note that the matrix multiplication in (12) evaluates to zero; there is no net contribution to Δs from the ΛK and ΛK^* states in the closure limit. There are in fact many such "sub-cancellations" in the closure sum for Δs: for each fixed set of spatial quantum numbers in the intermediate state, the sum over quark spins alone gives zero (because $\langle S_z^{(s)}\rangle = \langle S_z^{(\bar{s})}\rangle = 0$ in the ${}^3P_0$ state). That is, each $SU(6)$ multiplet inserted into Eq. (8) separately sums to zero.

Thus in this system, the closure sum to zero is enforced much more locally than in the $\omega - \phi$ mixing problem. After taking into account all energy denominators, the result of the calculation [10] is $\Delta s = -0.13$, in quite good agreement with the most recent extractions from experiment. It should be emphasized that all parameters of this calculation were fixed by spectra and decay data. Moreover, the result is quite stable to parameter changes.

The calculations of R_s^2 and μ_s are more difficult than the calculation of Δs, giving the results $R_s^2 = -0.04\text{fm}^2$ and $\mu_s = +0.035\ \mu_N$. For reasonable parameter variations, R_s^2 ranges between -0.02 and -0.06fm^2, but μ_s is quite stable.

Note on the Spin Crisis

With this background in mind, let me make some comments on the spin crisis. In the spirit of "valence quark plus glue with $q\bar{q}$ corrections", let us write

$$\Delta q = \Delta q_{valence} + \Delta q_{sea}$$

and note that:

1. Given the earlier discussion, we do not expect the nonrelativistic result $\Delta q_{valence} = 1$ since the lower components of the relativistic valence quarks developed *via* Z-graphs typically reduce their contributions to $\Delta q_{valence} \simeq 0.75$.

2. Since $\Delta q_{sea} = \sum_f \Delta q_{sea}^{(f)}$, where $\Delta q_{sea}^{(f)}$ is the spin sum contribution of the quark-antiquark sea of flavor f, if there are N_f approximately flavor-symmetric light quark flavors then $\Delta q_{sea} \simeq N_f \Delta q_{sea}^{(f_1)}$, where f_1 is the first of these light flavors. Note that no matter how suppressed $\Delta q_{sea}^{(f_1)}$ might be, if $N_f >> N_c$, $\Delta q - \Delta q_{valence}$ will be large. In other words, although the spin crisis makes the valence approximation look bad, what is relevant for the approximation is that $\Delta q_{sea}^{(s)} << \Delta q_{valence}$ which is indeed what is observed experimentally.

3. A possible scenario for the spin crisis is that $\Delta q_{valence} \simeq 0.75$, $\Delta q_{sea}^{(s)} \simeq -0.13$, $\Delta q_{sea}^{(u)} \simeq -0.16$, and $\Delta q_{sea}^{(d)} \simeq -0.16$ (where we have speculatively included a small $SU(3)$-breaking effect) leading to $\Delta q \simeq 0.3$. If this scenario is correct, then the spin crisis will have shown us [16] that the valence quarks behave just as they were supposed to do!

We can expect that, within the intrinsic systematic errors, Δu, Δd,and Δs will be known in another year or two. Then, the next logical step will be to determine the contribution of sea quarks, and the strange quarks in particular, to the static properties of the nucleons. Using parity violation as a probe, the SAMPLE experiment at MIT's Bates Lab and an extensive program of measurements planned for CEBAF at Jefferson Lab (including measurements utilizing the existing Hall A spectrometers as well as a new special purpose detector called G^0 funded for construction in Hall C) will allow us to decompose the nucleon form factors into their quark-level components: G_E^u, G_E^d, G_E^s, and G_M^u, G_M^d, G_M^s each as a function of Q^2.

Nonresonant Processes and $q\bar{q}$ Contributions to Charge Radii

It is natural to associate the valence quark model with resonance dominance of dynamical processes, and correspondingly to associate the $q\bar{q}$ content of hadrons with nonresonant effects. While there is an element of truth in these associations, they are also misleading. Consider, for example, the weak decay of a b quark in the $\bar{B}^0$ meson (I take this example because the current $b \to c$ is a *nonsinglet current* which cannot act on the sea quarks). Now take an extreme picture in which the $\bar{B}^0$ is full of $q\bar{q}$ pairs, but they are all totally frozen in place along a rigid flux tube connecting the b quark to a light valence $\bar{d}$ quark (an extreme adiabatic limit). With these dynamics, the $q\bar{q}$ sea would simply "go along for the ride" in the weak decay, which would be dominated by $c\bar{d}$ resonances, and their presence would not affect the analog of the charge radii for these weak transitions [17].

The cases we consider in baryon spectroscopy and dynamics are more complex. However, because in the closure approximation the net current on the $q\bar{q}$ pair is zero, our currents to some extent behave like non-singlet currents, so to the extent the sea is localized on the flux tube, it will also not produce nonresonant final states nor contribute to charge radii.

5 Key Issues for the Future

Let me close with a survey of some key issues I hope to see resolved by Baryon 2001.

5.1 Confirm Three Valence Degrees of Freedom in Baryons

Fundamental to our whole perspective on baryon spectroscopy, as well as our belief that this potentially very complex system simplifies, is the picture that the low-lying spectrum is indeed dominated by the degrees of freedom associated with three spin one-half quarks and their two relative coordinates.

We are now on the threshold of being able to test this belief. Of course it *cannot* be demonstrated using the 0 $\hbar\omega$ and 1 $\hbar\omega$ states, since to excite both the $\vec{\rho}$ and $\vec{\lambda}$ relative coordinates simultaneously one needs at least 2 $\hbar\omega$. We also note that the 2 $\hbar\omega$ states that are well-established are best interpreted as $\ell_\rho = 2, \ell_\lambda = 0$ and $\ell_\rho = 0, \ell_\lambda = 2$ states, which also obviously cannot demonstrate fully (and certainly not simultaneously) the expected degrees of freedom. This is of course the essence of the famous "missing resonance" problem: the ground state and 1 $\hbar\omega$ states cannot speak to the true numbers of degrees of freedom, and with only 8/21 of the 2 $\hbar\omega$ states established, there is no definitive evidence for the full anticipated spectrum of the valence quark model. We can all hope that by Baryon 2001 this problem will have been resolved. Bernhard Mecking will show you some early data from the CLAS spectrometer which adds fuel to this hope.

5.2 Check That Flux Tube Dynamics is the Relevant Gluodynamics

As described above, lattice QCD is supporting the flux tube picture. Recent results include detailed studies of the hybrid adiabatic surfaces as well as direct

calculations of the masses of the light hybrid mesons. Unfortunately, experiment isn't cooperating: recent results from BNL provide evidence for many $J^{PC} = 1^{-+}$ exotic signals with masses and properties that are at odds with theoretical expectations. These conflicts need to be resolved.

5.3 Determine the Origin of Spin-Dependent Forces

It had until recently been accepted dogma that spin-dependent forces arise in hadrons from one gluon exchange (OGE). A number of of papers [18] have recently suggested that one pion exchange (OPE) (or more generally, the exchange of the octet of pseudo-Goldstone bosons) and not OGE dominate spin-dependent interactions in baryons.

These papers legitimately point to a number of problems with the OGE model [19]:

- The OGE model does not explain why spin-orbit forces observed in baryons are small,
- The low mass of the 2 $\hbar\omega$ Roper resonance is "unnatural" in the OGE model, lying as it does below the 1 $\hbar\omega$ states,
- The $\Lambda(1520) - \Lambda(1405)$ splitting is not explained by the standard *ansatz* where spin-orbit forces are simply discarded,

while the OPE-based "chiral quark model" claims to solve some of these problems. In particular, it is noted by the proponents that

- OPE produces no spin-orbit forces, and
- Since the OPE hyperfine interaction is spin and flavor-dependent, it can be arranged to produce a pattern of splittings in which the Roper resonance is low-lying.

I have many objections to these proposals. My first objections are of a basic nature. As described above, the introduction of $q\bar{q}$ pairs into the quark model must be done with considerable care if its spectroscopy and the OZI rule are not to be destroyed. In doing so one will certainly discover that OPE, in the form of meson loop graphs that start on one quark and terminate on another, will play a role. The problem is stopping there, since the duality and associated closure properties I described above require that *all* meson loops be considered, and guarantees that higher mass meson loops will tend to cancel the OPE loops.

However, instead of dwelling on theoretical issues which might be turned into aesthetic ones, let me instead list what I consider to be some very practical problems with the proposed dynamics:

1. It is *not* an advantage that OPE produces *zero* spin orbit forces, since in such a situation strong spin-orbit forces are guaranteed! This is because we know on general grounds that the confining forces produce Thomas precession, which will make strongly-split *inverted* spin-orbit multiplets unless cancelled by some dynamical spin-orbit terms. In mesons the OGE and confining Thomas precession do seem to cancel nicely, and indeed as realized by early workers in baryons, such a cancellation also takes place in baryons for the analogous two-body forces [19]. Unfortunately, in baryons there are intrinsic three-body terms which no one to date has understood how to cancel.

2. The OGE mechanism is universal: the same basic process appears to control the hyperfine splittings in both mesons and baryons, and in both light quark and heavy quark systems. *E.g.*, that $\rho - \pi$ is twice $\Delta - N$ is natural to OGE. This model universality is also consistent with basic principles from heavy quark symmetry where hyperfine effects arise from the strength of the color-magnetic field produced by the light quarks at the location of the heavy quark. The OPE mechanism is in contrast very uneconomical: it could in principle work in baryons made of light quarks, but it cannot be responsible for the hyperfine interactions in light- or heavy-quark mesons, nor in heavy-quark baryons.

3. The successes of the OGE model are not just spectroscopic. Indeed, one of its earliest successes was the prediction of the mixing angles required in the negative parity N^*'s to produce the observed pattern of strong decays. In particular, the strong mixing in the $N^*\frac{1}{2}^-$ sector, explaining the large $N\eta$ branching ratio of the $N^*(1535)$, was an early success of the OGE models [19]. These well-established mixing angles are not reproduced by the OPE model.

4. I have already alluded to the constraints of heavy quark symmetry. The OPE model is not relevant to the physics of the heavy quark limit. In contrast, the OGE model automatically incorporates the constraints and theorems of heavy quark symmetry in $Q\bar{q}$ and Qqq systems when one quark mass m_Q is much larger than Λ_{QCD} and perturbative QCD in $Q_1\bar{Q}_2$ and $Q_1Q_2Q_3$ when all quark masses are large (recall Figs. 1 and 2). The OGE model in fact does an excellent job in explaining the hyperfine

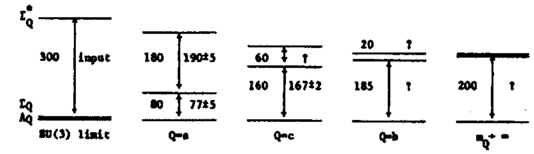

Figure 6: The udQ baryon spectra as a function of the "heavy" quark mass.

interactions in $\bar{B}$ and D mesons, and *with the same parameters* describes the $\bar{K}$ and π systems. As mentioned just above, in this case OGE also seems to do a reasonable job of explaining the spin-orbit splittings [20]. In the heavy quark limit of udQ baryons, it is a theorem that the Σ_Q^* and Σ_Q are degenerate, with the Λ_Q below them. In the OGE model, this $(\Sigma_Q^*, \Sigma_Q) - \Lambda_Q$ splitting is predicted to be about two-thirds of the observed $\Delta - N$ splitting [21]. This is qualitatively what is seen in charmed baryons. In fact the OGE model explains quantitatively the hyperfine interactions in udc baryons, including the splittings of Λ_c, Σ_c, and Σ_c^*. Moreover, once again *with the same parameters*, it describes the uds baryons Λ, Σ, and Σ^* and in the light quark sector the N and Δ (see Fig. 6; note that since the publication of this Figure, the $\Sigma_c^* - \Sigma_c$ splitting has been measured to be 65 MeV, close to what is shown). This makes it difficult to believe that the physics of the $SU(3)$ baryons and mesons is qualitatively different from the physics of heavy baryons.

5. Finally, the $\Lambda(1520) - \Lambda(1405)$ splitting, which is held up as a failure of the OGE model, is not readily explained by the OPE model, where it has been suggested that it be interpreted as a $\bar{K}N$ bound state. In fact I suspect that this system is a place to look for a solution of the OGE spin orbit problem: recent data on excited charmed baryons has made it very likely that these two states are the s-quark analogs of a heavy-quark multiplet with $s_\ell^{\pi_\ell} = 1^-$, as they are supposed to be in the OGE-based valence quark model. If verified, this would rule out a $\bar{K}N$ interpretation of the $\Lambda(1405)$ and enforce the standard quark model interpretation. Indeed, these states are predicted [19] to have a very simple

internal structure in which they consist of a spin-zero ud diquark in an $L = 1$ orbital excitation relative to the heavy quark Q or s. This internal structure has long been supported by a mixing angle analysis of the decay patterns of the $\Lambda^* \frac{1}{2}^-$ baryons, but the new charmed baryon data lends great credence to this prediction. It also suggests that this system might be amenable to a meson-like analysis of spin-orbit forces to the extent that the ud diquark behaves like a single object orbiting the heavy quark Q[22].

5.4 Explore the Nature of the Scalar Mesons

There are several influences converging that should focus our attention on the scalar meson sector over the next few years:

- It is now very likely, based on lattice QCD and modelling, that the primitive 0^{++} glueball is in the 1.5 GeV mass region. Moreover, it is likely to be the only low mass glueball: others are predicted to lie above 2 GeV, and, in a blow to any hopes that at least some glueballs could be easily extracted from conventional $q\bar{q}$ signals, there are no low mass exotic glueballs expected.
- Experimentally, glueball hunting is heating up: there are new and convincing sightings of scalar states in the relevant mass region.
- One test of the proposal presented above for the survival of the OZI rule is the prediction that it should fail dramatically in the 0^{++} sector: in that unique case the closure sum need not zero since the created and destroyed $q\bar{q}$ pair has vacuum quantum numbers. We should therefore expect the scalar meson sector to behave very differently from any other meson nonet we have encountered. *I.e.,* no miracles will happen here, and warnings are hereby issued to glueball hunters to beware of the scalar meson sector.

Unravelling the physics of this sector will not be made easier by the fact that the two most obvious states with scalar quantum numbers, the $f_0(980)$ and the $a_0(980)$, are widely interpreted as candidates for $qq\bar{q}\bar{q}$ states, and in particular as "$K\bar{K}$ molecules".

From the preceeding one might correctly conclude that the scalar meson sector is interesting in its own right. However, for me its primary interest lies in the fact that it contains vital clues to the low energy structure of QCD. It seems quite possible that the connection between the current quarks and gluons of QCD and the valence quarks and potentials of the quark model

is in the (necessarily scalar) $q\bar{q}$ vacuum condensate. Among other things, these structures, which exist in the vacuum because of confinement even in the quenched approximation, can give the massless QCD quarks their constituent quark masses.

5.5 Determine the Spectrum of Multiquark Sectors

We should never forget that one of the main motivations for studying hadrons is to eventually be able to place our understanding of nuclear physics on a firmer foundation. Some of the most profound questions in strong interaction physics are questions of this type, including

- Why can the nucleus be described in the "nuclear physics approximation" as being made of A nucleons instead of a soup of $3A$ quarks?
- What is the origin of the residual forces between these nucleonic clusters? and
- What is the nature of nuclear matter in the short-distance regime where these clusters necessarily overlap and lose their identity as low energy effective degrees of freedom?

At a very concrete level, we can approach these problems through better experimental and theoretical understanding of such systems as the H dibaryon (is it a true six quark state, a nuclear physics-type state, or is it a continuum system consisting of multiple channels $\Lambda\Lambda$, $N\Xi$, $\Sigma\Sigma$, *etc.*?). I have already mentioned two other concrete systems which need to be understood: the scalar mesons $f_0(980)$ and $a_0(980)$, which seem to be nuclear physics-type states ("molecules") and the $\Lambda(1405)$, which has been a strong candidate for a $\bar{K}N$ state, but now seems more likely to be an ordinary uds baryon.

6 Summary

In this introductory talk I have advocated approaching the phenomenology of QCD in two steps: in zeroth order, strong QCD should be approximated by a relativistic constituent quark model with flux tube gluodynamics; as a second step, $q\bar{q}$ sea and other $1/N_c$ effects should be consistently added as perturbations.

I have described how the second step, "unquenching" the quark model, might be done in a way that preserves its spectroscopic successes and respects the OZI rule. All of the results reported are qualitative, but the approach appears to be a viable route to explaining the underlying physics. The key assumptions of the model are that flux tube dynamics (including a flux tube full of $q\bar{q}$ pairs) can be treated adiabatically, and that the $q\bar{q}$ pair creation occurs into a state with vacuum quantum numbers ($^3P_0$ or $J^{PC} = 0^{++}$).

The main new results of the picture I have advocated are:

1. The quark model spectrum is immune to meson-loop-induced mass shifts apart from those associated with nearby thresholds. By systematically incorporating the adiabatic effects into the definition of the quark model potential, a systematic low energy expansion of the effects of thresholds is possible [11].

2. The OZI rule survives loops corrections, once they are done systematically; an exception to this rule will occur for $J^{PC} = 0^{++}$ mesons since their mixing is not zero in the closure limit [9].

3. Both of the preceeding results are associated with insights into the hadronic states required to respect the dualities which underlie them, and they strongly suggest that low energy hadronizations of QCD are in trouble, since sums over large towers of states were required.

4. $\Delta s \simeq -0.13$, which when combined with comparably polarized u and d sea quarks and with $\Sigma_{valence} \sim 0.75$, suggests [16] that the valence quarks are actually "normal", with the sea quarks responsible for the spin crisis.

With this picture and these new results for guidance, I have also outlined some of the key isues that I hope to hear discussed here at Baryons 98 and resolved between now and Baryon 2001. These include:

- confirm there are three valence degrees of freedom in baryons,
- check that flux tube dynamics is the relevant gluodynamics,
- determine the origin of spin-dependent forces,
- explore the nature of the scalar mesons both to find the glueball and as a window on the $q\bar{q}$ vacuum condensate, and
- determine the spectrum in multiquark sectors such as $qqqq\bar{q}$, $qqq\bar{q}\bar{q}$, and $qqqqqq$.

Since I hope to see at least some of these matters resolved here, let me stop immediately so that we can officially begin Baryons 98!

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